UPR-1198-T

A string theoretic model of gauge mediated supersymmetry beaking

Mirjam Cvetič and Timo Weigand
Department of Physics and Astronomy, University of Pennsylvania, Philadelphia, USA

We propose a robust supergravity model of dynamical supersymmetry breaking and gauge mediation, and a natural embedding in non-perturbative string theory with D-branes. A chiral field (and its mirror) charged under "anomalous" U(1)'s acts as a Polonyi field whose hierarchical Polonyi-term can be generated by string instantons. Further quartic superpotential terms arise naturally as a tree-level decoupling effect of massive string states. A robust supersymmetry breaking minimum allows for gauge mediation with soft masses at the TeV scale, which we realise for a globally consistent SU(5) GUT model of Type I string theory, with a D1-instanton inducing the Polonyi term.

I. Introduction One of the mysteries of particle physics is if Nature has chosen supersymmetry to protect the electroweak scale in the TeV regime and, if so, how the breaking of supersymmetry is communicated to the Standard Model. Among the different mediation mechanisms, gauge mediation [1] offers a clean rationale for the absence of flavour-changing neutral currents. Its model-independent experimental signatures [2, 3] make it an interesting scenario also in light of future testability in the Large Hadron Collider (LHC) era. In its simplest implementation, supersymmetry is broken by the F-term of a hidden gauge sector chiral superfield S and communicated to the visible sector by a vector pair of messenger fields q, $\tilde{q}$. Gaugino and slepton masses arise at the one- and two-loop level, respectively, roughly of order $\frac{\alpha}{4\pi}\frac{F}{S}$, while the gravitino mass is of order F/M_{Pl} (M_{pl}-Planck mass). One of the challenges in realising gauge mediation is therefore to dynamically generate a sufficiently small vacuum expectation value (VEV) for the supersymmetry breaking field $\langle S\rangle < 10^{-3}\, M_{Pl}$ in order for the effects of gravity mediation to be subleading [15].

The perhaps simplest supersymmetry breaking scenario involves a linear superpotential of Polonyi type $\mu^2 S$ whose F-term breaks supersymmetry at the scale μ^2. In string theory, D-brane instantons [5, 6, 7] can account both for the presence of the Polonyi term [8] and for a hierarchical suppression of its scale μ, as demonstrated even in globally consistent examples [9]. Yet, in order to realise gauge mediation, further dynamical input is required to stabilise S at the desired hierarchical scale. This can be achieved [10] by including one-loop Kähler potential corrections with a specific sign [16].

In this article we propose an alternative model where the Polonyi field is stabilised with the help of quartic superpotential terms. These generically result from integrating out heavy string states provided the Polonyi field is charged under massive $U(1)$ gauge factor(s). The specific terms we are interested in require a vector pair $\tilde{S}$ with opposite $U(1)$ charge(s). Models of this type arise naturally in string theory with D-branes. Under specific assumptions on the moduli dynamics this framework possesses a supersymmetry breaking vacuum tailor-made for gauge mediation to generate TeV scale soft masses. Within Type I compactifications with D-instantons we also present a globally consistent Grand Unified Theory (GUT) model of the type discussed in [9] where the stringy consistency conditions allow for TeV soft masses.

II. The model Our supersymmetry breaking hidden sector consists of a massive $U(1)_a \times U(1)_b$ gauge theory with bi-fundamental chiral superfields $S_{(-1_a,1_b)}$ and $\tilde{S}_{(1_a,-1_b)}$ and

$$W = \mu^2 S + c - \frac{S^2\tilde{S}^2}{4M} - \lambda S q\tilde{q}, \tag{1}$$

$$K = SS^\dagger + \tilde{S}\tilde{S}^\dagger + qq^\dagger + \tilde{q}\tilde{q}^\dagger. \tag{2}$$

The messenger fields q_{-1_b}, $\tilde{q}_{1_a}$ form a vector-like pair under the visible sector gauge group and will be responsible for gauge mediation of supersymmetry breaking. As will be detailed in the next section the above model is well-motivated from string theory: The massless vector-like pair S and $\tilde{S}$ arises for instance at the intersection of a pair of D6-branes in the Type IIA context. The gauge bosons of $U(1)_a$ and $U(1)_b$ acquire string scale masses via the Green-Schwarz mechanism. The perturbatively forbidden Polonyi term $\mu^2 S$ for a charged field can naturally be generated by D-brane instantons [9][17]. Gauge invariance is ensured by the shift of the axionic part of the closed modulus T which determines the part of the instanton action charged under the massive $U(1)$s [5, 6, 7] in $\mu^2 = \mu_0^2\, e^{-T}$. Higher monomials in S or $\tilde{S}$ only are perturbatively absent due to charge selection rules. The constant c results from integrating out the closed string moduli uncharged under $U(1)_a$ and $U(1)_b$, which are assumed to be stabilised at a scale much higher than μ.

The non-renormalisable quartic term in the superpotential arises from [14] the integration of massive (closed or suitable open string sector) modes C with mass M_C which couple in the superpotential to $S\tilde{S}$ as

$$W_C = \lambda_C\, CS\tilde{S} + M_C\, C^2\,, \tag{3}$$

with $M = M_C/\lambda_C^2$. For heavy closed string states and $\lambda_C \sim 1$, $M \sim M_s$ – the string scale. By contrast, if M_C is associated with dynamically generated masses for (closed and open sector) string moduli, M could be $\ll M_s$. The effective M can in principle decrease significantly due to enhanced threshold effects of higher mass level string states C whose multiplicity increases exponentially. On

the other hand if there were a selection rule that would set $\lambda_C = 0$, the effective $M \to \infty$ and the massive states would decouple. We shall see that our results are extremely robust in that they depend only mildly on the values of M and are basically driven by $\mu \ll 1$ in the interplay between the Polonyi and the tree-level quartic superpotential term.

The superpotential W_C (3) also induces tree-level Kähler potential corrections [14] $\delta K_1 = +\frac{SS^\dagger \tilde{S}\tilde{S}^\dagger}{4\,M^2}$. K can furthermore receive one-loop corrections due to the couplings of S and $\tilde{S}$ to bi-linears of open-string heavy states with mass $\tilde{M}$, $\delta K_2 \sim -\frac{(SS^\dagger)^2}{\Lambda^2} - \frac{(\tilde{S}\tilde{S}^\dagger)^2}{\tilde{\Lambda}^2}$, where $\Lambda \sim \tilde{\Lambda} \sim \pi \tilde{M}$. As we shall see momentarily, for our explicit solution $S \sim \tilde{S} \ll 1$ both δK_1 and δK_2 can be neglected relative to the superpotential term $-\frac{S^2 \tilde{S}^2}{4M}$.

The full scalar potential $V = V_F + V_D$ takes the standard $\mathcal{N} = 1$ supergravity form

$$\begin{aligned} V_F &= e^K (D_i W\, D_{\bar{j}} W K^{i\bar{j}} - 3|W|^2), \\ V_D &= \tfrac{g_a^2}{2}(-|S|^2 + |\tilde{S}|^2 + |q|^2 + \xi_a)^2 + \\ &\quad \tfrac{g_b^2}{2}(|S|^2 - |\tilde{S}|^2 - |\tilde{q}|^2 + \xi_b)^2 \end{aligned} \tag{4}$$

in terms of $D_i W = \partial_i W + K_{,i}\, W$ and the gauge couplings $g_{a,b}$ associated with $U(1)_a$ and $U(1)_b$, respectively. We take $M_{Pl} = 1$. The Fayet-Iliopoulos (FI) terms ξ_a and ξ_b depend on T and in general also on other closed moduli N_i not entering the superpotential (1). Their full dynamics hinges upon the precise form of their Kähler potential [13]. To avoid such model dependent questions we do not analyse the stabilisation of T and N_i explicitly here but assume they can be stabilised such that the FI terms are at most of order $\mu^{8/3}$ (see later); furthermore under specific conditions on the Kähler potential for T and N_i their backreaction on the effective potential (4) is negligible. An analysis of these conditions will be presented elsewhere.

Under these assumptions on the moduli sector there exists a stable solution in the regime $\langle S\rangle \sim \langle \tilde{S}\rangle = \mathcal{O}(\mu^2\, M)^{\frac{1}{3}}) \ll 1$ and $\langle q\rangle = \langle \tilde{q}\rangle = 0$, as ensured by their F-terms. We reparametrise S and $\tilde{S}$ as $S = |S|\exp(i\phi)$, $\tilde{S} = |\tilde{S}|\exp(i\tilde{\phi})$ and proceed iteratively by first enforcing vanishing D-terms via $\langle |S|\rangle = \langle |\tilde{S}|\rangle \equiv s$, neglecting the FI terms at this stage. It is convenient to introduce a new combination of fields $S_\pm = (|S| \pm |\tilde{S}|)/\sqrt{2}$, where S_- obtains the dominant mass $m^2_{S_-} = 4(g_a^2 + g_b^2)s^2$ from the D-term. The supergravity potential V_F in (5) is in this approximation expanded only in terms of $(S_+,\, \phi,\ \tilde{\phi})$. Since we are looking for a minimum in the region $\langle S_+\rangle = \mathcal{O}((\mu^2 M)^{\frac{1}{3}}) \ll 1$ it suffices to expand only up to terms proportional to $\mu^4 \langle S_+\rangle = \mu^4\sqrt{2}s$. The leading order potential then takes the form

$$V_{F_0} = \mu^4 - 3c^2 - \frac{\mu^2 s^3}{M}\cos(\sqrt{5}\phi_1) + \frac{s^6}{2M^2}\,, \tag{5}$$

where we have introduced new fields $\phi_1 \equiv (\phi + 2\tilde{\phi})/\sqrt{5}$, and $\phi_2 \equiv (-2\phi + \tilde{\phi})/\sqrt{5}$. Note that up to the term $-3c^2$ this is the potential obtained in the globally supersymmetric approximation. The minimum of (5) and its zero value there are ensured by taking respectively

$$\phi_1 = 0\,, \quad s = \left(\mu^2 M\right)^{\frac{1}{3}}\ ; \quad c = \mu^2/\sqrt{6}\,, \tag{6}$$

while ϕ_2 has a flat direction. In the next step we correct for having set $\langle S_-\rangle = 0$ in V_{F_0} and find $\langle S_-\rangle = \mathcal{O}(\mu^2/((g_a^2 + g_b^2)\,M)) \ll \langle S_+\rangle$. The tiny deviation from $S_- = 0$ entails a subleading D-term of $\mathcal{O}(\mu^{\frac{8}{3}}/M^{\frac{2}{3}}) \ll F$. The backreaction from the closed string moduli dynamics is likewise subleading under the above assumption of FI terms scaling at most like $\mu^{8/3}$. This justifies the iterative procedure and the correction to $\langle S_+\rangle$ is negligible. At this order the potential respects R-symmetry and ϕ_2 is the Goldstone boson of this spontaneously broken global symmetry. The degeneracy of ϕ_2 is removed due to supergravity effects in the potential at the next order, linear in powers of s (for simplicity we set $\phi_1 = 0$ at this order),

$$V_{F_1} = -\frac{cs\mu^2}{2}(8 + \frac{s^3}{\mu^2 M})\cos(\frac{2\phi_2}{\sqrt{5}})\,, \tag{7}$$

which fixes $\phi_2 = 0$. At this order c and s are likewise corrected, but the corrections are suppressed by $\mathcal{O}(s) \ll 1$ and thus again subleading. The scalar S_- is much heavier than S_+ and ϕ_1, while ϕ_2 has a positive mass-square for positive c (at this level) and is further suppressed by one power of s For the values (6) one obtains

$$m^2_{S_-} = 4(g_a^2 + g_b^2)M^{\frac{2}{3}}\mu^{\frac{4}{3}}\,, m^2_{S_+} = \tfrac{9}{4}\, M^{-\frac{2}{3}}\mu^{\frac{8}{3}} = \tfrac{9}{10}m^2_{\phi_1}\,,$$

and

$$m^2_{\phi_2} = \tfrac{9}{5}\, c\, M^{-\frac{1}{3}}\mu^{\frac{4}{3}} = \tfrac{3\sqrt{6}}{10}M^{-\frac{1}{3}}\mu^{\frac{10}{3}}\,. \tag{8}$$

The model predicts $3/\sqrt{10}$ for the mass ratio of S_+ and ϕ_1. Note that only $m^2_{\phi_2}$ depends on the value of c. The mass-square correction $\delta_D m^2_{S+}$ due to the small D-term is of $\mathcal{O}(\mu^4/M^2)$ and thus subleading.

The F-term messenger masses are of the order λs and positive as long as $\mu^2 \le \lambda s^2$ or equivalently $\mu \le \lambda^{\frac{3}{2}} M$. This is satisfied for the relevant range of $\mu \sim 10^{-10}$ (see below) and $\lambda \gg 10^{-6}$. The D-term mass corrections $\delta_D m_q = \mathcal{O}(\lambda s\, \mu^{\frac{2}{3}}/M^{\frac{2}{3}})$ are again subleading. The model has a small gravitino mass $m_{3/2} = \mu^2$.

Coleman-Weinberg one-loop corrections due to the superpotential coupling of S to the messengers q and $\tilde{q}$ in (2) result in the Kähler potential correction

$$\delta K = -\kappa SS^\dagger \log(\tfrac{SS^\dagger}{\Lambda^2})\,, \qquad \kappa \equiv \frac{\lambda^2 N_c}{16\pi^2}\,, \tag{9}$$

where the renormalisation scale Λ at which the coupling λ is defined is chosen to be of the order of the VEV of S. N_c is the number of colors associated with the observable sector gauge group $SU(N_c)$, e.g., $N_c = 5$ for $SU(5)$ GUT, and the messengers are in respective $\mathbf{N_c}$ and $\overline{\mathbf{N}}_\mathbf{c}$ representations. Since these corrections respect R symmetry

they cannot modify the mass of ϕ_2 and their contribution to the masses of other scalars are only subleading for $\lambda \leq 1$ [18].

Phenomenological Analysis Due to the relative suppression of the D- versus the F-term, the supersymmetry soft masses are dominated by gauge-mediated F-term breaking. The loop-generated visible sector soft masses are determined by $m_{soft} \sim \frac{\alpha}{4\pi} \langle F \rangle / \langle S \rangle \sim 10^{-3}\mu^2/s$ [1] and lie in the TeV range provided $\mu^2 \sim 10^{-13}s$; in this case the solution (6) for s implies a relationship $\mu \sim 10^{-10}\ M^{\frac{1}{4}}$ and consequently $s \sim 10^{-7}M^{\frac{1}{2}}$. The corresponding hidden sector scalar masses are predicted to be in the range

$$m_{S_-} \sim 10^{10} - 10^{11}\,\text{GeV}, \qquad m_{S_+} \sim 10^3 - 10^4\,\text{TeV},$$
$$m_{\phi_1} \sim 10^3 - 10^4\,\text{TeV}, \qquad m_{\phi_2} \sim 1 - 10^2\,\text{TeV}\,.$$

The messenger masses are in the range $10^{10} - 10^{11}$ GeV. Interestingly, the model has a light gravitino of mass in the range $0.1 - 10$ GeV. In table I we present numerical values for a wider parameter range of μ and M.

μ	M	m_{S_-}	m_{S_+}	m_{ϕ_2}	m_q	$m_{3/2}$
10^9	10^{18}	$2.83\,10^{11}$	$1.50\,10^6$	$2.71\,10^2$	$1.00\,10^{11}$	0.10
10^9	10^{16}	$6.08\,10^{10}$	$6.96\,10^6$	$5.84\,10^2$	$2.15\,10^{10}$	0.10
10^{10}	10^{18}	$1.31\,10^{12}$	$3.23\,10^7$	$1.26\,10^4$	$4.64\,10^{11}$	10.0
10^{10}	10^{16}	$2.83\,10^{11}$	$1.50\,10^8$	$2.71\,10^4$	$1.00\,10^{11}$	10.0

TABLE I: Masses for $S_\pm$ and ϕ_2, messengers q and the gravitino for different values of μ and M. Note, $m_{\phi_1} = \sqrt{10}m_{S_+}/3$ and $m_{\tilde{q}} = m_q$. We took $\lambda = 0.10$, $g_a = g_b = 0.10$, $N_c = 5$ and $M_{Pl} = 1.22\,10^{19}$ GeV. All masses are in GeV.

The above analytic results are extremely close to actual numerical values of the potential: corrections to these analytic expressions, which would appear in the potential at the $\mu^4 s^2$ order, modify the above expressions at a level of $100 \times \mathcal{O}(s)\,\% \sim 10^{-5}\%$. Note also that Kähler potential corrections due to massive modes, as described above, contribute only at this order. Therefore even if we lower the string scale M and $\tilde{\Lambda}$ to, say, 10^{-3} and 10^{-2}, respectively, these corrections are small.

Other Minima There exists no nearby supersymmetric minimum with non-zero VEV for the messenger fields, unless there are additional fields $(q'_{1_b}, \tilde{q}'_{-1_a})$ to ensure D-flatness as $s \to 0$ and $\langle q \rangle = \langle \tilde{q} \rangle \to \mu/\sqrt{\lambda}$. This would lead to a supersymmetric vacuum with vacuum energy $-\mu^4/2$. Our solution is stable against false vacuum decay into this supersymmetric one as the bounce action can be estimated to be $\sim \pi^2(\Delta s)^4/\Delta V \sim \pi^2\mu^{8/3}/\mu^4 \sim 10^{30}$. In absence of mirror fields $(q', \tilde{q}')$, there exists only a non-supersymmetric nearby solution with $\langle S \rangle = \langle q \rangle = \langle \tilde{q} \rangle = \mu/\sqrt{3\lambda}$ and $\langle \tilde{S} \rangle = 0$ and positive energy $+\mu^4/6$. Note also that as $s \to \mathcal{O}(1)$ the Kähler potential corrections, discussed above, lead to a singular Kähler metric and the blow-up of the potential.

III. Global string realisation The described model of gauge mediation has a natural realisation in string theory. For definiteness our discussion focuses on intersecting D-brane models in Type IIA or Type IIB Calabi-Yau orientifolds where massless charged matter arises from strings stretching between different D-branes. Generalisations to the respective strong coupling M- or F-theory versions are likewise possible.

A particularly economic realisation appears in the context of $SU(5)$ GUT models. The hidden sector consists of two stacks of single D-branes wrapping (possibly magnetised) cycles Π_a, Π_b with a massless vector pair S and $\tilde{S}$ in the (a,b) sector. The SU(5) gauge group can arise from a stack of 5 coincident branes on Π_c with the $\overline{\mathbf{5}}_m$ and $[\mathbf{5}_H + \overline{\mathbf{5}}_H]$ localised at intersections with another single brane Π_d. In addition, the setup contains the orientifold image of each brane stack and matter in the, say, (a,b) sector is identified with the image sector and (b', a'). A possible identification of the visible, hidden and messenger sector matter is given in table II.

Particle	Charge	Sector	Particle	Charge	Sector
(Q_L, U^c_R, e^c_R)	$\mathbf{10}_{2c}$	(c', c)	S	$(-1_a, 1_b)$	(a,b)
(L, D^c_R)	$\overline{\mathbf{5}}_{(1_d,-1_c)}$	(c,d)	$\tilde{S}$	$(1_a, -1_b)$	(b,a)
Higgs	$\mathbf{5}_{(1_d,1_c)}$	(d', c)	q	$\mathbf{5}_{(-1_b,1_c)}$	(b,c)
N^c_R	1_{-2d}	(d,d')	$\tilde{q}$	$\overline{\mathbf{5}}_{(1_a,-1_c)}$	(c,a)

TABLE II: Embedding of supersymmetry breaking hidden sector into $U(5)_c \times U(1)_d$ GUT theory.

For the sake of applications we now specialise to compactifications of Type I string theory on a Calabi-Yau manifold X. The relevant D-branes are space-filling D9-branes carrying holomorphic vector bundles V_a. Their orientifold image carries the dual bundle $V_a^\vee$. We will choose the simple case of line bundles L_a. The massless open matter in the, say, (a,b), sector is counted by the cohomology group $H^i(X, L_a^\vee \otimes L_b)$, where $i=1$ refers to chiral matter in the bi-fundamental $(\overline{N}_a, N_b)$ and $i=2$ to the conjugate representation $(N_a, \overline{N}_b)$. For technical details see [9].

The Polonyi term $\mu^2 S$ in the superpotential can be generated by Euclidean D1-branes, so-called E1-instantons, which are localised in the four external dimensions and wrap suitable holomorphic 2-cycles C on X. For this to happen the instanton has to carry precisely one charged fermionic zero mode λ_a and λ_b of charge $+1_b$ and -1_b, respectively. These arise from massless open strings between the instanton and the two hidden brane stacks with bundles L_a and L_b and are counted by $H^i(C, L_a|_C \otimes \sqrt{K_C})$. Here $i=0$ and $i=1$ refer to chiral modes of charge 1_a and -1_a, respectively [19]. The scale of the resulting superpotential is set by the string coupling g_s and the volume of the instanton cycle Vol_C in string units as $W_{np} = M_s^2 \exp(-\frac{2\pi}{g_s}\text{Vol}_C)$.

We now demonstrate in an explicit, globally consistent GUT toy model how the scales leading to TeV soft masses in the visible sector can arise. Our main assumption is that the quartic superpotential term is due to massive

Bundle	N	$c_1(L) = q\sigma + \pi^*(\zeta)$
L_a	1	$\pi^*(-l + 2E_1 + 2E_2 - 2E_3 - E_4)$
L_b	1	$4\sigma + \pi^*(l - 2E_2 + E_4)$
L_c	5	$\pi^*(2E_1 - 2E_2 - 2E_3)$

TABLE III: A $U(5)_c \times U(1)_a \times U(1)_b$ Polonyi model.

string modes of string scale M_s such that $M = \mathcal{O}(M_s)$. In Type I theory, the string scale is determined by the string coupling g_s as $M_s^2 = 2\pi\, m_g^2\, g_s\, \alpha_{GUT}$ with $m_g \simeq 2.4 \times 10^{18}$ GeV. If we assume that g_s is stabilised in the perturbative regime, say, $g_s = 0.4$, then $\alpha_{GUT} = 0.04$ implies $M_s = 7.6 \times 10^{17}$ GeV and Vol_C has to be stabilised around 2.63 for TeV soft masses.

While the actual stabilisation of g_s and the geometric moduli is not addressed in this paper, we now show in our concrete example that the above sample values are compatible with the non-trivial D-term supersymmetry constraints that arise from the presence of the magnetised D9-branes. As in [9] we work on an elliptic fibration over dP_4. The SU(5) GUT brane Π_c and the hidden sector Π_a, Π_b correspond to D9-branes with line bundles L_c, L_a and L_b and multiplicities N as in table III. For details of the notation see [9]. The D5-brane tadpole is cancelled by introducing D5-branes along curves in the effective class $W = 24F + \pi^*(16l - 4E_1 - 12E_2 - 4E_3)$, and the K-theory charge $\sum_i N_i c_1(L_i)$ mod 2 vanishes. The toy model gives rise to 4 chiral families of $\mathbf{10}$ plus additional vector pairs. In contrast to the general setup of table II there is no $U(1)_d$ stack, but the $\overline{\mathbf{5}}_m$ and $[\mathbf{5}_H + \overline{\mathbf{5}}_H]$ arise from the sector between Π_c and the D5-branes. From $H^*(L_a^* \otimes L_b) = (0, 27, 5, 0)$ one reads off a considerable excess of hidden sector fields in addition to the minimal $(S, \tilde{S})$ pair as in the scenario of the previous section.

The Polonyi term for S arises from an instanton wrapping an isolated $\mathbb{P}^1$ in the divisor $\pi^* E_4$ with zero intersection with the base of the fibration [9]. Indeed we find $H^*(C, L_a|_C \otimes \sqrt{K_C}) = (1, 0)$ and $H^*(C, L_b|_C \otimes \sqrt{K_C}) = (0, 1)$ and thus precisely the amount of charged zero modes λ_a and λ_b required to generate the Polonyi term. As the instanton cycle is rigid, no subtleties associated with extra uncharged zero modes arise. The D-term supersymmetry conditions can be satisfied inside the Kähler cone, e.g., for $r_\sigma = 1.06$, $r_l = 9.37$, $r_1 = -4.99$, $r_2 = r_3 = -3.00$, $r_4 = -2.63$. The resulting gauge kinetic function $f = \tilde{f}/(2\pi g_s)$ with $\tilde{f}_a = 12$ leads to a slightly too small value of $\alpha_{GUT} = g_s/\tilde{f}_a = 0.03$. The instanton volume is given by $|r_4| = 2.63$, as required for $\mu^2 = \mathcal{O}(10^{-20})$ and TeV soft masses. While the spectrum of the visible sector is by no means semi-realistic, the example does serve as a prototype demonstrating how our model of gauge mediated supersymmetry breaking can be engineered in string theory. We hope to implement this module into more realistic string vacua.

Acknowledgements We thank G. Shiu for collaboration in initial stages and P. Ouyang and G. Shiu for many important discussions. We acknowledge useful discussions with J. Marsano, Y. Nomura and R. Richter. We thank the Aspen Center for Physics for hospitality during the course of this research. T.W. thanks the University of Madison and the Banff Center for hospitality. Our work is supported in part by the DOE Grant DOE-EY-76-02-3071, the NSF RTG DMS Grant 0636606 and the Fay R. and Eugene L. Langberg Endowed Chair.

[1] G. F. Giudice and R. Rattazzi, Phys. Rept. **322** (1999) 419 [arXiv:hep-ph/9801271].
[2] P. Meade, N. Seiberg and D. Shih, arXiv:0801.3278 [hep-ph].
[3] J. Distler and D. Robbins, arXiv:0807.2006 [hep-ph].
[4] D. E. Diaconescu, B. Florea, S. Kachru and P. Svrcek, JHEP **0602** (2006) 020 [arXiv:hep-th/0512170].
[5] R. Blumenhagen, M. Cvetič and T. Weigand, Nucl. Phys. B **771** (2007) 113 [arXiv:hep-th/0609191].
[6] L. E. Ibáñez and A. M. Uranga, JHEP **0703** (2007) 052 [arXiv:hep-th/0609213].
[7] B. Florea, S. Kachru, J. McGreevy and N. Saulina, JHEP **0705** (2007) 024 [arXiv:hep-th/0610003].
[8] O. Aharony, S. Kachru and E. Silverstein, Phys. Rev. D **76**, 126009 (2007) [arXiv:0708.0493 [hep-th]].
[9] M. Cvetič and T. Weigand, Phys. Rev. Lett. **100**, 251601 (2008) [arXiv:0711.0209 [hep-th]].
[10] R. Kitano, Phys. Lett. B **641** (2006) 203 [arXiv:hep-ph/0607090].
[11] H. Murayama and Y. Nomura, Phys. Rev. D **75** (2007) 095011 [arXiv:hep-ph/0701231].
[12] M. Buican and S. Franco, arXiv:0806.1964 [hep-th].
[13] M. Cicoli, J. P. Conlon and F. Quevedo, arXiv:0805.1029 [hep-th].
[14] M. Cvetič, L. L. Everett and J. Wang, Nucl. Phys. B **538**, 52 (1999) [arXiv:hep-ph/9807321].
[15] For an early discussion within string theory see [4].
[16] For a recent alternative approach see e.g. [11], and for a stringy mediation mechanism using D-instantons [12].
[17] The same instanton may also generate the terms of the type $S^2\tilde{S}$ whose coupling is of the order of μ^2/M_s^2 (M_s-string scale) and thus for small values of S and $\tilde{S}$ it is negligible compared to the Polonyi term. We thank E. Dudas for a discussion on this point.
[18] Other one-loop corrections due to the (S, $\tilde{S}$) sector lead to Kähler potential corrections of the type $\sim \frac{(SS^\dagger)^2}{M^2}\log(\frac{SS^\dagger}{\Lambda^2})$ and are thus subleading. A deviation from D-flatness, due to a small non-zero S_-, leads to negligible one-loop corrections proportional to $S_-^2 S_+^2$.
[19] Note the slight change in conventions as compared to [9].